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Resonant H and A Mixing in CP–noninvariant 2HDM and MSSM

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In the general two–Higgs–doublet model, including the minimal supersymmetric standard model (MSSM) as its specific realization, the two heavy neutral Higgs bosons are nearly degenerate in the decoupling limit. If the theory is CP–noninvariant, the mixing between the heavy states can strongly be affected by the decay widths. We develop the formalism describing this CP–violating non–hermitian mixing and provide some interesting experimental signatures of the CP–violating mixing at a $\gamma\gamma$ collider with polarized beams in the context of the CP–noninvariant MSSM.

1. INTRODUCTION

The MSSM is a specific realization of general scenarios that include two doublet fields in the Higgs sector. After three Goldstone fields are absorbed by electroweak gauge bosons, the remaining five fields give rise to physical states. At the tree level, the MSSM Higgs sector is CP invariant, with two CP–even and one CP–odd neutral states. However, the MSSM offers new sources of CP violation, which render the Higgs sector CP–noninvariant at the loop level. In such CP–noninvariant theories the three neutral states mix to form a triplet with both even and odd components in the wave–functions under CP transformations [1, 2, 3]. The mixing can become very large if the states are nearly mass–degenerate. This situation is naturally realized for supersymmetric theories in the decoupling limit [4] in which two of the neutral states are heavy.

In the present report we describe a simple quantum mechanical (QM) formalism for the CP–violating resonant H/A mixing in the decoupling limit and then discuss some experimental signatures of the CP–violating mixing in Higgs production and decay processes at a photon collider with polarized photon beams.

2. MIXING FORMALISM

The self-interaction of two Higgs doublets in a CP–noninvariant theory is generally described by the potential [4]

$$\begin{aligned}\mathcal{V} \;=\; & m_{11}^2\Phi_1^\dagger\Phi_1 + m_{22}^2\Phi_2^\dagger\Phi_2 - [m_{12}^2\Phi_1^\dagger\Phi_2 + \mathrm{h.c.}] + \frac{1}{2}\lambda_1(\Phi_1^\dagger\Phi_1)^2 + \frac{1}{2}\lambda_2(\Phi_2^\dagger\Phi_2)^2 \\ & +\lambda_3(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) + \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) + \left\{\frac{1}{2}\lambda_5(\Phi_1^\dagger\Phi_2)^2 + \left[\lambda_6(\Phi_1^\dagger\Phi_1) + \lambda_7(\Phi_2^\dagger\Phi_2)\right]\Phi_1^\dagger\Phi_2 + \mathrm{h.c.}\right\}\end{aligned} \tag{1}$$

where $\Phi_{1,2}$ denote two complex $Y = 1$, $\mathrm{SU}(2)_L$ iso-doublet scalar fields. The coefficients are in general all non–zero. The parameters $m_{12}^2, \lambda_{5,6,7}$ can be complex, incorporating the CP-noninvariant elements in the interactions. The neutral components of the scalar fields Φ_1 and Φ_2 are assumed to develop non–zero vacuum expectation values (vevs) $\langle\phi_1^0\rangle = v_1/\sqrt{2}$ and $\langle\phi_2^0\rangle = v_2/\sqrt{2}$, which can be chosen real and positive without loss of generality. As usual, $v = (v_1^2 + v_2^2)^{1/2} = 246$ GeV.

It is useful to rotate the Higgs fields $\Phi_{1,2}$ to the $\Phi_{a,b}$ basis with the angle β satisfying $\tan\beta = v_2/v_1$ (we exploit the abbreviations $t_\beta = \tan\beta$, $c_\beta = \cos\beta$, $s_{2\beta} = \sin 2\beta$ etc.). In this basis only the field Φ_a develops a non–zero vev

$$\Phi_a = \begin{pmatrix} G^+ \\ \frac{1}{\sqrt{2}}\left(v + H_a + iG^0\right)\end{pmatrix}, \qquad \Phi_b = \begin{pmatrix} H^+ \\ \frac{1}{\sqrt{2}}\left(H_b + iA\right)\end{pmatrix} \tag{2}$$

and the three fields $G^{\pm,0}$ can be identified as the would-be Goldstone bosons, while $H^{\pm}, H_{a,b}$ and A give rise to physical Higgs bosons. The real mass matrix $\mathcal{M}_0^2$ of neutral Higgs fields in the basis of H_a, H_b, A, which is hermitian and symmetric by CPT invariance, can easily be derived from the Higgs potential (1) after the rotations:

$$\mathcal{M}_{0R}^2 = v^2 \begin{pmatrix} \lambda & -\hat{\lambda} & -\hat{\lambda}_p \\ -\hat{\lambda} & \lambda - \lambda_A + M_A^2/v^2 & -\lambda_p \\ -\hat{\lambda}_p & -\lambda_p & M_A^2/v^2 \end{pmatrix} \tag{3}$$

after eliminating $m_{11,22}^2, m_{12}^{2I}$ from the minimization conditions, and exchanging m_{12}^{2R} for the *auxiliary* parameter M_A^2. It is defined by the relation

$$m_{12}^{2R} = \frac{1}{2}(M_A^2 s_{2\beta} + v^2(\lambda_5^R s_{2\beta} + \lambda_6^R c_\beta^2 + \lambda_7^R s_\beta^2)) \tag{4}$$

and it will be one of the key parameters in the system. The $\lambda, \hat{\lambda}$ and λ_A parameters are functions of the real parts, while λ_p and $\hat{\lambda}_p$ are functions of the imaginary parts of the parameters λ_i in Eq. (1); their explicit form can be found in Ref. [5].

In a CP–invariant theory all λ_i couplings are real and the off–diagonal elements $\lambda_p, \hat{\lambda}_p$ vanish. Thus the neutral mass matrix breaks into the CP–even 2×2 part, and the [stand–alone] CP–odd part. The 2×2 part gives rise to two CP–even neutral mass eigenstates h and H, while M_A is identified as the mass of the CP–odd Higgs boson A. In the CP–violating case, however, all three states mix leading to $H_{1,2,3}$ mass eigenstates with no definite CP parities.

For small mass differences, the mixing of the states is strongly affected by their widths. This is a well–known phenomenon for resonant mixing [6] and has also been recognized for the Higgs sector [7]. The hermitian mass matrix (3) has therefore to be supplemented by the anti–hermitian part $-iM\Gamma$ incorporating the decay matrix [8]

$$\mathcal{M}^2 = \mathcal{M}_0^2 - iM\Gamma \tag{5}$$

This matrix includes the widths of the Higgs states in the diagonal elements as well as the transition elements within any combination of pairs. They are particularly important in the case of nearly mass–degenerate states. All these elements $(M\Gamma)_{ab}^{AB}$ are built up by loops of the fields (AB) in the self-energy matrix $\langle h_a h_b \rangle$ of the Higgs fields.

In general, the light Higgs boson, the fermions and electroweak gauge bosons, and in supersymmetric theories, gauginos, higgsinos and scalar states may contribute to the loops in the propagator matrix. In the physically interesting case of decoupling, the mixing structure simplifies considerably allowing for a very simple and transparent approach [5]. Alternative approach requires a full coupled–channel analysis [9].

2.1. Decoupling Limit

The decoupling limit [4] is defined by the inequality $M_A^2 \gg |\lambda_i|\, v^2$ with $|\lambda_i| \lesssim O(1)$. In this limit the H_a state becomes the CP–even light Higgs boson h and decouples from H_b and A. The heavy states $H = H_b$ and A are nearly mass degenerate, which turns out to be crucial for large mixing effects between H and A. It is therefore enough to consider a lower–right 2×2 submatrix of the matrix (3) for the heavy H/A states which we write as follows

$$\mathcal{M}_{HA}^2 = \begin{pmatrix} M_H^2 - iM_H\Gamma_H & \Delta_{HA}^2 \\ \Delta_{HA}^2 & M_A^2 - iM_A\Gamma_A \end{pmatrix} \tag{6}$$

where Δ_{HA}^2 also consists of a real dissipative part and an imaginary absorptive part. Moreover, the couplings of the heavy Higgs bosons to gauge bosons and their supersymmetric partners are suppressed. In the case of all supersymmetric particle contributions to be suppressed either by couplings or by phase space in $M\Gamma$, it is sufficient to consider only loops built up by the light Higgs boson and top quark; for the explicit form of the light Higgs boson and top quark loop contributions to the matrix $M\Gamma$, we refer to Ref. [5]. The loops also contribute to the real part of the mass matrix, either renormalizing the λ parameters of the Higgs potential or generating such parameters if not present yet at the tree level.

2.2. Physical Masses and States

The symmetric complex mass–squared matrix $\mathcal{M}^2$ in Eq.(6) can be diagonalized through a *complex rotation*

$$\mathcal{M}^2_{H_iH_j} = \begin{pmatrix} M^2_{H_2} - iM_{H_2}\Gamma_{H_2} & 0 \\ 0 & M^2_{H_3} - iM_{H_3}\Gamma_{H_3} \end{pmatrix} = C\mathcal{M}^2_{HA}C^{-1} \tag{7}$$

where the mixing matrix and the mixing angle are given by

$$C = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix}, \qquad X = \frac{1}{2}\tan 2\theta = \frac{\Delta^2_{HA}}{M^2_H - M^2_A - i\left[M_H\Gamma_H - M_A\Gamma_A\right]} \tag{8}$$

A non–vanishing (complex) mixing parameter $X \neq 0$ requires CP–violating transitions between H and A either in the real mass matrix, $\lambda_p \neq 0$, or in the decay mass matrix, $(M\Gamma)_{HA} \neq 0$, [or both]. However, note that even for nearly degenerate masses, the mixing could be suppressed if the widths were significantly different. As a result, the mixing phenomena are strongly affected by the form of the decay matrix $M\Gamma$. Since the difference of the widths enters through the denominator in X, the modulus $|X|$ becomes large for small differences and small widths.

The mixing shifts the Higgs masses and widths in a characteristic pattern [6]. The two complex mass values after and before diagonalization are related by the complex mixing angle θ:

$$M^2_{H_3} - M^2_{H_2} - i\left(M_{H_3}\Gamma_{H_3} - M_{H_2}\Gamma_{H_2}\right) = \left[M^2_A - M^2_H - i(M_A\Gamma - M_H\Gamma_H)\right] \times \sqrt{1+4X^2} \tag{9}$$

Since the eigenstates of the complex, non–hermitian matrix $\mathcal{M}^2$ are no longer orthogonal, the ket and bra mass eigenstates have to be defined separately: $|H_i\rangle = C_{i\alpha}|H_\alpha\rangle$ and $\langle\widetilde{H}_i| = C_{i\alpha}\langle H_\alpha|$ ($i = 2, 3$ and $H_\alpha = H, A$). The final state F in heavy Higgs formation from the initial state I is, then, described with the amplitude

$$\langle F|H|I\rangle = \sum_{i=2,3} \langle F|H_i\rangle \frac{1}{s - M^2_{H_i} + iM_{H_i}\Gamma_{H_i}} \langle\widetilde{H}_i|I\rangle \tag{10}$$

where the sum runs only over diagonal transitions in the mass–eigenstate basis.

3. EXPERIMENTAL SIGNATURES

To illustrate the general QM results in a realistic example, we adopt a specific MSSM scenario with the source of CP–violation localized in the complex trilinear coupling A_t of the soft supersymmetry breaking part involving the top squark.[1] All other interactions are assumed to be CP–conserving. For $\phi_A \neq 0, \pi$, the stop–loop corrections induce the CP–violation in the effective Higgs potential (1). The effective λ_i parameters have been calculated in Ref. [3] to two–loop accuracy; to illustrate the crucial points we take the dominant one–loop $t/\tilde{t}$ contributions.

More specifically, we take a typical set of parameters from Ref. [11],

$$M_S = 0.5 \text{ TeV}, \quad |A_t| = 1.0 \text{ TeV}, \quad \mu = 1.0 \text{ TeV}; \quad \tan\beta = 5 \tag{11}$$

and change the phase ϕ_A of the trilinear parameter A_t. With $\phi_A = 0$ we find the following values of the light and heavy Higgs masses and decay widths, and the stop masses:

$$M_h = 129.6\,\text{GeV},\ M_H = 500.3\,\text{GeV},\ M_A = 500.0\,\text{GeV};\ \ \Gamma_H = 1.2\,\text{GeV},\ \Gamma_A = 1.5\,\text{GeV};\ \ m_{\tilde{t}_{1/2}} = 372/647\,\text{GeV} \tag{12}$$

Clearly, with the mass splitting of 0.3 GeV, the heavy Higgs states are not distinguishable. When the phase ϕ_A is

[1]This assignment is compatible with the bounds on CP–violating SUSY phases from experiments on electric dipole moments [10].

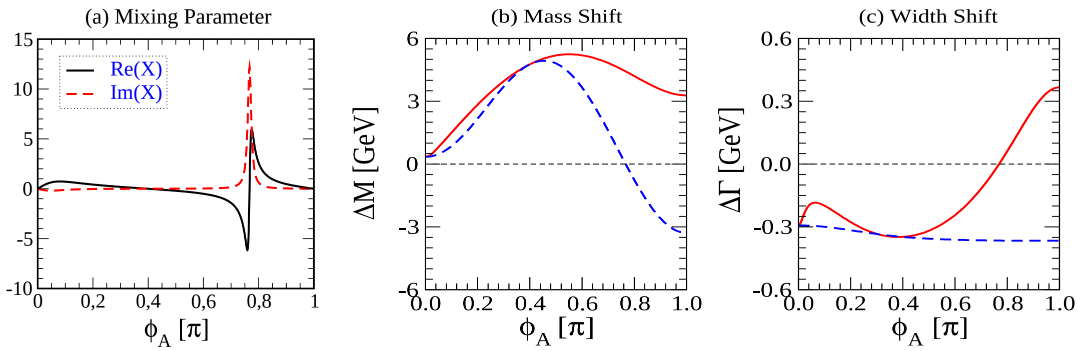

Figure 1: The ϕ_A dependence of (a) the mixing parameter X and of the shifts of (b) masses and (c) widths with the phase ϕ_A evolving from 0 to π for $\tan\beta = 5$, $M_A = 0.5$ TeV and couplings as specified in the text; in (b,c) the mass and width differences without mixing are shown by the broken lines. $\Re\mathrm{e}/\Im\mathrm{m}X(2\pi - \phi_A) = +\Re\mathrm{e}/-\Im\mathrm{m}X(\phi_A)$ for angles above π.

turned on,[2] the CP composition, the masses and the decay widths of heavy states are strongly affected, as shown in Figs. 1(a), (b) and (c), while the mass of the light Higgs boson h is not. The heavy two–state system shows a very sharp resonant CP–violating mixing, purely imaginary a little above $\phi_A = 3\pi/4$, Fig. 1(a). The mass shift is enhanced by more than an order of magnitude if the CP–violating phase rises to non-zero values, reaching a maximal value of ~ 5.3 GeV; the width shift changes between -0.3 and $+0.4$ GeV. As a result, the two mass–eigenstates should become clearly distinguishable at future colliders, in particular at a photon collider [12]. Moreover, both states have significant admixtures of CP–even and CP–odd components in the wave–functions. Since $\gamma\gamma$ colliders offer unique conditions for probing the CP–mixing [13, 14, 15], we discuss two experimental examples: (a) Higgs formation in polarized $\gamma\gamma$ collisions and (b) polarization of top quarks in Higgs decays, where spectacular signatures of resonant mixing can be expected.

(a) The amplitude of the reaction $\gamma\gamma \to H_i \to F$ is a superposition of H_2 and H_3 exchanges. For equal helicities $\lambda = \pm 1$ of the two photons, the amplitude reads

$$\mathcal{M}^F_\lambda = \sum_{i=2,3} \langle F|H_i\rangle \frac{1}{s - M^2_{H_i} + iM_{H_i}\Gamma_{H_i}} \left[S^\gamma_i(s) + i\lambda P^\gamma_i(s)\right] \tag{13}$$

where $\sqrt{s}$ is the $\gamma\gamma$ energy and the loop–induced $\gamma\gamma H_i$ scalar and pseudoscalar form factors, $S^\gamma_i(s)$ and $P^\gamma_i(s)$, are related to the well–known conventional $\gamma\gamma H/A$ form factors, $S^\gamma_{H,A}$ and $P^\gamma_{H,A}$. For their relation and explicit form we refer to Refs. [5] and [11]. In our scenario the Higgs–tt couplings are assumed to be CP–conserving, implying negligible top-loop contributions to P^γ_H and S^γ_A since the gluino mass is sufficiently heavy compared with the stop masses, while the $\tilde{t}_1$ loop generates a non-negligible CP-violating amplitude S^γ_A. In the region of strong mixing on which we focus, however, the CP-violating vertex corrections have only a small effect on the experimental asymmetries compared with the large impact of CP-violating Higgs-boson mixing.

Polarized photons provide a very powerful tool to investigate the CP properties of Higgs bosons. With linearly polarized photons one can project out the CP–even and CP–odd components of the H_i wave–functions by arranging the photon polarization vectors to be parallel or perpendicular. On the other hand, circular polarization provides us with a direct insight into the CP-violating nature of Higgs bosons. Two asymmetries are of interest

$$\mathcal{A}_{lin} = \frac{\sigma_\parallel - \sigma_\perp}{\sigma_\parallel + \sigma_\perp}, \qquad \mathcal{A}_{hel} = \frac{\sigma_{++} - \sigma_{--}}{\sigma_{++} + \sigma_{--}} \tag{14}$$

where $\sigma_\parallel$, $\sigma_\perp$ and σ_{++}, σ_{--} are the corresponding total $\gamma\gamma$ fusion cross sections for linear and circular polarizations, respectively. Though CP–even, the asymmetry $\mathcal{A}_{lin}$ can serve as a powerful tool nevertheless to probe CP–violating

[2] With one phase ϕ_A, the complex mixing parameter X obeys the relation $X(2\pi - \phi_A) = X^*(\phi_A)$, implying all CP–even quantities symmetric and all CP–odd quantities anti–symmetric about π.

admixtures to the Higgs states since $|\mathcal{A}_{lin}| < 1$ requires both S_i^γ and P_i^γ non-zero couplings. A more direct probe of CP–violation due to H/A mixing is provided by the CP–odd (and also CP$\tilde{\text{T}}$–odd) asymmetry $\mathcal{A}_{lin}$.

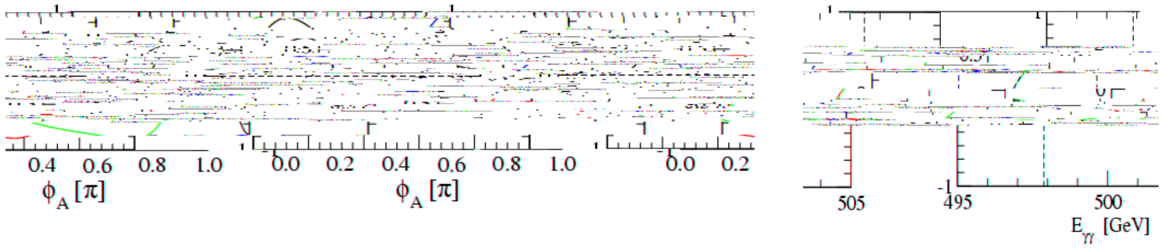

Figure 2: The ϕ_A dependence of the CP–even and CP–odd correlators, $\mathcal{A}_{lin}$ (left panel) and $\mathcal{A}_{hel}$ (center panel), at the poles of H_2 and H_3, respectively, and the $\gamma\gamma$ energy dependence (right panel) of the correlators, $\mathcal{A}_{lin,hel}$, for $\phi_A = 3\pi/4$ in the production process $\gamma\gamma \to H_i$ in the limit in which H/A mixing is the dominant CP–violating effect. The same parameter set as in Fig. 1 is employed. The vertical lines on the right panel mark positions of the two mass eigenvalues, M_{H_3} and M_{H_2}.

Fig. 2 show the ϕ_A dependence of the asymmetries $\mathcal{A}_{lin}$ and $\mathcal{A}_{hel}$ at the poles of H_2 and of H_3, respectively, for the same parameter set as in Fig. 1 and with the common SUSY scale $M_{\tilde{Q}_3} = M_{\tilde{t}_R} = M_S = 0.5$ TeV for the soft SUSY breaking top squark mass parameters. By varying the $\gamma\gamma$ energy from below M_{H_3} to above M_{H_2}, the asymmetries, $\mathcal{A}_{lin}$ (blue solid line) and $\mathcal{A}_{hel}$ (red dashed line), vary from -0.39 to 0.34 and from -0.29 to 0.59, respectively, as demonstrated on the right panel of Fig. 2 with $\phi_A = 3\pi/4$, a phase value close to resonant CP–mixing.

(b) A second observable of interest is the polarization of the top quarks in H_i decays produced by $\gamma\gamma$ fusion or elsewhere in various production processes at an e^+e^- linear collider and LHC $H_{2,3} \to t\bar{t}$. Even if the H/Att couplings are [approximately] CP–conserving, the complex rotation matrix C may mix the CP–even H and CP–odd A states leading to CP–violation. In the production–decay process $\gamma\gamma \to H_i \to t\bar{t}$, two CP–even and CP–odd correlators between the transverse t and $\bar{t}$ polarization vectors $s_\perp, \bar{s}_\perp$

$$\mathcal{C}_\parallel = \langle s_\perp \cdot \bar{s}_\perp \rangle \qquad \text{and,} \qquad \mathcal{C}_\perp = \langle \hat{p}_t \cdot (s_\perp \times \bar{s}_\perp) \rangle \tag{15}$$

can be extracted from the azimuthal–angle correlation between the two decay planes $t \to bW^+$ and $\bar{t} \to \bar{b}W^-$ [13].

Fig. 3 shows the ϕ_A dependence of the CP–even and CP–odd asymmetries, $\mathcal{C}_\parallel$ and $\mathcal{C}_\perp$, at the poles of H_2 and of H_3, left and center panels respectively. If the invariant $t\bar{t}$ energy is varied throughout the resonance region, the correlators $\mathcal{C}_\parallel$ (blue solid line) and $\mathcal{C}_\perp$ (red dashed line) vary characteristically from -0.43 to -0.27 [non–uniformly] and from 0.84 to -0.94, respectively, as shown in the right panel of Fig. 3.

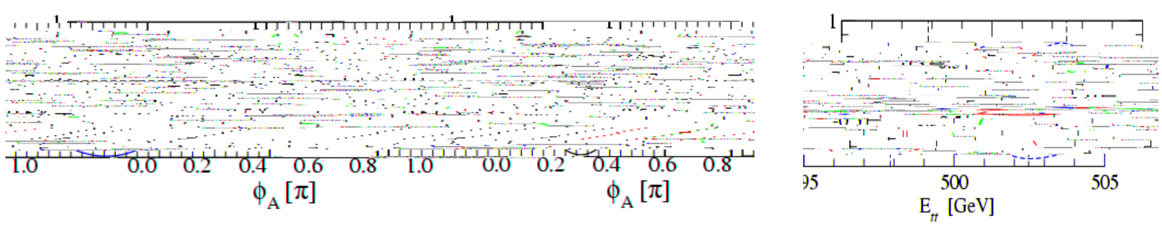

Figure 3: The ϕ_A dependence of the CP–even and CP–odd correlators, $\mathcal{C}_\parallel$ (left panel) and $\mathcal{C}_\perp$ (center panel), at the pole of H_2 and H_3 and the invariant $t\bar{t}$ energy dependence (right panel) of the correlators $\mathcal{C}_{\parallel,\perp}$ for $\phi_A = 3\pi/4$ in the production–decay chain $\gamma\gamma \to H_i \to t\bar{t}$. [Same SUSY parameter set as in Fig.2.]

4. CONCLUSIONS

Exciting mixing effects can occur in the supersymmetric Higgs sector if CP–noninvariant interactions are present. In the decoupling regime these effects can become very large, leading to interesting experimental consequences. Higgs formation in $\gamma\gamma$ collisions with polarized beams proves particularly interesting for observing such effects. However, exciting experimental effects are also predicted in such scenarios for $t\bar{t}$ final–state analyses in decays of the heavy Higgs bosons at LHC and in the e^+e^- mode of linear colliders.

Detailed experimental simulations would be needed to estimate the accuracy with which the asymmetries presented here can be measured. Though not easy to measure, the large magnitude and the rapid, significant variation of the CP–even and CP–odd asymmetries through the resonance region with respect to both the phase ϕ_A and the $\gamma\gamma$ energy would be a very interesting effect to observe in any case.

Acknowledgments

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